Heavy chiral bosons search at hadron colliders

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The production of new spin-1 chiral bosons at the hadron colliders, the Fermilab Tevatron and the CERN LHC, is considered. The masses of the chiral bosons can be determined on the basis of experimental data of precise low-energy experiments, which already indicate indirectly their existence. They can explain, for example, the serious 4.5σ discrepancy between the measured and the predicted two pion branching ratio of the τ decay and the sign of the 3.3σ deviation of the muon (g-2) theoretical prediction from the experimental value. Quantitative evaluations of the various differential cross-sections of the chiral boson production at hadron colliders are made using the CalcHEP package. It is noteworthy that the Tevatron data already hint the existence of the lightest charged chiral boson with a mass around 500 GeV. New Tevatron data and the LHC results will definitely confirm or reject this indication. In the positive case the LHC would be able to discover all predicted charged and neutral chiral bosons spanning in mass up to 1 TeV.

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I. INTRODUCTION

The hadron colliders, due to their biggest center-of-mass energy and their relatively compact sizes, still remain a main tool for discoveries of very heavy particles. So in 1983 the two dedicated UA1 [1] and UA2 [2] experiments discovered the intermediate vector bosons at the CERN SPS Collider. The collider energy, 540 GeV, just met the condition for a rough estimation of the minimal center-of-mass energy about six times the predicted mass of the weak bosons. The factor six corresponds to the fact that the only part of the proton momentum is shared by the (anti)quarks. At such energy the production cross-sections of both W and Z bosons have the nanobarn level. One faces, however, a very large background from the strong interactions.

In order to detect the production of the heavy bosons, only events with bosons pure leptonic decays into isolated high transverse-momentum leptons and without prominent associated jet activity have been selected. Such selection supplies backgroundless conditions for the detection of the resonantly produced charged W bosons and more than three orders of magnitude smaller background from the Drell–Yan dilepton production under the neutral Z boson peak. At the same time the signal from the decays of these bosons will play the role of a background for searching of new heavy intermediate bosons, W' and Z', with similar couplings to the quarks and leptons.

In any case, besides the simple manifestation of existence of the weak bosons, one needs precise study of their properties following from the Standard Model (SM). This task has been excellently fulfilled by the Large Electron-Positron (LEP) storage ring at CERN and the Stanford Linear Collider (SLC) at SLAC. Unfortunately, the masses of the t quark and the undiscovered yet Higgs boson happened to be too high to be discovered at these colliders. Nevertheless, the precision of the electroweak measurements at the lepton colliders was so high, that the predicted from radiative loop corrections mass of the top-quark $m_t = 180\,^{+8}_{-9}\,^{+17}_{-20}$ GeV [3] has been found in agreement and with comparable accuracy of its first direct measurements at the Fermilab Tevatron by the CDF [4] $m_t = 176 \pm 8 \pm 10$ GeV and the D0 [5] $m_t = 199\,^{+19}_{-21} \pm 22$ GeV collaborations.

In spite of the overwhelming background for the top-quark pair production in the strong interactions at the hadron collider, the uncertainty of the top-quark mass $m_t = 170.9 \pm 1.1 \pm 1.5$ GeV [6] is considerably reduced at present. Moreover, recently, the evidence for a single top-quark production [7] through the weak interactions and direct measurement of $|V_{tb}|$ at the Fermilab Tevatron hadron collider became possible. Another achievement in precise measurements at the hadron collider is the W-mass measurement $m_W = 80.413 \pm 0.048$ GeV [8] by the CDF collaboration at comparable with the LEP experiments accuracy, which represents the single most precise measurement to date. All these measurements will allow further constrain the mass of the Higgs particle, which discovery is the main priority task of the running Tevatron and the constructing Large Hadron Collider (LHC).

Discovery of the predicted theoretically heavy particles and establishing of the SM without any surprises are signifying the experimental high energy physics for the last thirty years. Therefore, the LHC construction is connected not only with the Higgs discovery, but with the hope to find the physics beyond the SM. Needless to say, that the discovery of new particles and new interactions at the hadron colliders requires a strong input of theoretical models in order to extract the signals over huge background.

The aim of the present paper is to point out the signature for the resonant production of heavy spin-1 *chiral* bosons and their decays. It has become proverbial (see, for example, the textbook [9]), that the *Jacobian peak* in the transverse momentum/mass distribution is characteristic of all two-body decays. It is not, however, the case for the decay of the new chiral bosons [10]. To my knowledge such bosons for the first time were introduced

by Kemmer [11] and are naturally appeared in the extended conformal supergravity theories [12].

In the next section I give the motivation of introduction of such kind new bosons and describe their model-independent properties. Further in section III I suggest a simple model for the extension of the SM in order to make more definite predictions. Its phenomenological consequences are used for explanation of the anomalies in the precise low-energy experiments and as guiding line for the detection and identification of such bosons at the hadron colliders, the Fermilab Tevatron and the CERN LHC. In section IV various distributions are produced using the CalcHEP package [13] and the first indications of the production of the lightest charged chiral bosons at the Tevatron are presented. In conclusion the prospects for the LHC physics and the modification of the PYTHIA program [14] for simulations of the chiral boson production and their decays are discussed.

II. INTRODUCTION OF THE CHIRAL BOSONS AND THEIR MODEL-INDEPENDENT PROPERTIES

The trilinear interactions, wherein various bilinear fermion currents couple to appropriate bosons with dimensionless coupling constants, play a key role in the interactions of the elementary particles. They include both the gauge couplings of the (axial-)vector fermion currents to the (axial-)vector bosons and the Yukawa couplings of the (pseudo)scalar fermion currents to the (pseudo)scalar bosons. Up to now only these interactions have found phenomenological applications. For example, they have been used as effective couplings of various mesons to quarks/baryons in the model description of the strong nuclear interactions. On a more fundamental level of the elementary particles, the gauge interactions and the Yukawa couplings of the Higgs bosons to the matter constitute the base of the SM. The latter seem today a viable solution of acquiring masses of the initially massless fermions.

In this paper I consider additional trilinear interactions, which naturally appear in the full set of derivativeless couplings of fermion currents to bosons and which have been missed in phenomenological applications. Let us consider all bilinear hermitian combinations of the fermion fields. In the relativistic case spin-$\frac{1}{2}$ fermion fields are described by the *four*-component Dirac *bispinor* ψ. Therefore, there are 16 independent bilinear combinations $\bar\psi\mathcal{O}\psi$:

$$\begin{aligned} \mathcal{S} = \bar\psi\psi, \qquad \mathcal{P} = i\bar\psi\gamma^5\psi, \\ \mathcal{V}^\mu = \bar\psi\gamma^\mu\psi, \qquad \mathcal{A}^\mu = \bar\psi\gamma^\mu\gamma^5\psi, \\ \mathcal{T}^{\mu\nu} = \bar\psi\sigma^{\mu\nu}\psi, \end{aligned} \tag{1}$$

which can couple to the S scalar, P pseudoscalar, V_μ vector, A_μ axial-vector and $T_{\mu\nu}$ rank-2 antisymmetric tensor fields, correspondingly,

$$\begin{aligned} \mathcal{L}_Y &= g_S\ \bar\psi\psi\ S + ig_P\ \bar\psi\gamma^5\psi\ P \\ &+ g_V\ \bar\psi\gamma^\mu\psi\ V_\mu + g_A\ \bar\psi\gamma^\mu\gamma^5\psi\ A_\mu \\ &+ \frac{t}{2}\ \bar\psi\sigma^{\mu\nu}\psi\ T_{\mu\nu}, \end{aligned} \tag{2}$$

with dimensionless coupling constants.

All these bilinear combinations represent either spin-0 or spin-1 states in accordance with the relation $\mathbf{\frac{1}{2}} \otimes \mathbf{\frac{1}{2}} = \mathbf{0} \oplus \mathbf{1}$. However, in contrast to the non-relativistic case, there are *two different* spin-$\frac{1}{2}$ fermions, left-handed $\psi_L = \frac{1}{2}(1-\gamma^5)\psi$ and right-handed $\psi_R = \frac{1}{2}(1+\gamma^5)\psi$, which transform under *inequivalent* representations of the Lorentz group ($\frac{1}{2}$,0) and (0,$\frac{1}{2}$), correspondingly. Therefore, the number of the independent combinations is twice enlarged in comparison with the non-relativistic case. The new scalar $\bar\psi\gamma^0\psi$ and pseudoscalar $\bar\psi\gamma^0\gamma^5\psi$ components are associated with the fourth components of the relativistic vector $\bar\psi\gamma^\mu\psi$ and axial-vector $\bar\psi\gamma^\mu\gamma^5\psi$ currents and do not lead to any new physical spin-0 states. While the new $\bar\psi\sigma^{0i}\psi$ vector and $\bar\psi\sigma^{ij}\psi$ axial-vector spin-1 states are introduced as components of the independent second rank antisymmetric tensor current $\bar\psi\sigma^{\mu\nu}\psi$.

In other words, in the relativistic case there are *two different* spin-1 states which transform under the *inequivalent* vector ($\frac{1}{2}$,$\frac{1}{2}$) and *chiral* (1,0)+(0,1) representations of the Lorentz group. The former four-component representation is associated with well-known (axial-)vector bosons, A_μ and V_μ, while the latter six-component representation demands introduction of a new rank-2 antisymmetric tensor field, $T_{\mu\nu}$, which describes simultaneously three-components vector and axial-vector bosons.

The new tensor Yukawa coupling in (2) can be rewritten explicitly in its chiral form

$$\mathcal{L}_Y^T = \frac{t}{2\sqrt{2}}\ \overline{\psi_L}\sigma^{\mu\nu}\psi_R\ T^+_{\mu\nu} + \frac{t}{2\sqrt{2}}\ \overline{\psi_R}\sigma^{\mu\nu}\psi_L\ T^-_{\mu\nu}, \tag{3}$$

where chiral components $T^\pm_{\mu\nu} = \frac{1}{\sqrt{2}}(T_{\mu\nu} \pm i\tilde{T}_{\mu\nu})$ of the antisymmetric tensor field and its dual $\tilde{T}_{\mu\nu} = \frac{1}{2}\epsilon_{\mu\nu\alpha\beta}T^{\alpha\beta}$ were introduced. These components are connected through the C charge-conjugate and the P parity transformations in the Minkowski space and are (anti)selfdual tensors in the euclidean space.

In order to find the free Lagrangian for the new rank-2 antisymmetric tensor field it is enough to evaluate the one-loop radiative correction into its self-energy from virtual fermion pairs with known propagators. Since the Yukawa coupling constant t is dimensionless, the theory is formally renormalizable and the structure of the quantum corrections should reproduce the bare free Lagrangian. Simple calculations lead to conformally invariant Lagrangian

$$\mathcal{L}_0^T = \frac{1}{4}\ \partial_\rho T_{\mu\nu}\ \partial^\rho T^{\mu\nu} - \partial_\mu T^{\mu\rho}\ \partial^\nu T_{\nu\rho} \tag{4}$$

for the antisymmetric tensor field with mass dimension one [15]. The properties of this Lagrangian have been investigated in [16]. Here I just mention that in contrast to the gauge bosons, the massless vector and axial-vector bosons, which are described by the antisymmetric tensor field, have only longitudinal physical components.

It is interesting to note that the mass term $T_{\mu\nu}T^{\mu\nu}$ is not generated by the quantum corrections, because it is protected by the chiral symmetry with transformations

$$\begin{aligned} \psi &\rightarrow \exp[i\theta\gamma^5]\psi, & \bar{\psi} &\rightarrow \bar{\psi}\exp[i\theta\gamma^5], \\ T^+_{\mu\nu} &\rightarrow \exp[-2i\theta]T^+_{\mu\nu}, & T^-_{\mu\nu} &\rightarrow \exp[+2i\theta]T^-_{\mu\nu}. \end{aligned} \tag{5}$$

In the same time the chirally invariant selfinteraction

$$\mathcal{L}^T_4 = \lambda\left(T_{\mu\nu}T^{\mu\nu}T_{\alpha\beta}T^{\alpha\beta} - 4T_{\mu\nu}T^{\nu\alpha}T_{\alpha\beta}T^{\beta\mu}\right). \tag{6}$$

is generated even in case of a *real* antisymmetric tensor field, due to its chiral properties [17].

The chiral symmetry can be localized $\theta \rightarrow \theta(x)$ including the axial-vector gauge fields with transformations

$$A_\mu \rightarrow A_\mu + \frac{1}{g_A}\,\partial_\mu\theta(x). \tag{7}$$

In this case additional gauge interactions of the antisymmetric tensor field are necessary [17]

$$\begin{aligned} \mathcal{L}^T_A &= 2g_A\left(\tilde{T}^{\mu\nu}\partial^\rho T_{\rho\nu} - T^{\mu\nu}\partial^\rho\tilde{T}_{\rho\nu}\right)A_\mu \\ &+ g_A^2\left(A_\rho A^\rho T_{\mu\nu}T^{\mu\nu} - 4A_\mu A^\nu T^{\mu\rho}T_{\nu\rho}\right). \end{aligned} \tag{8}$$

These interactions lead to the negative contribution in the β-function of the gauge coupling constant g_A, which exhibits *asymptotically free* behavior even in the *abelian* case.

The chiral bosons acquire the mass analogously to the gauge bosons through the Higgs mechanism and a *nonlocal* mass term is generated [18]. Using the relation [19]

$$T_{\mu\nu} = \hat{R}_{\mu\nu} - \frac{1}{2}\,\epsilon_{\mu\nu\alpha\beta}\,\hat{B}^{\alpha\beta}, \tag{9}$$

where $\hat{R}_{\mu\nu} = \hat{\partial}_\mu R_\nu - \hat{\partial}_\nu R_\mu$, $\hat{B}_{\mu\nu} = \hat{\partial}_\mu B_\nu - \hat{\partial}_\nu B_\mu$ and $\hat{\partial}_\mu = \partial_\mu/\sqrt{-\partial^2}$, the free Lagrangian (4) for the antisymmetric tensor field can be rewritten through the field-strength tensors of the new vector $R_\mu = \hat{\partial}^\nu T_{\nu\mu}$ and axial-vector $B_\mu = \hat{\partial}^\nu \tilde{T}_{\nu\mu}$ fields in a more convenient form

$$\mathcal{L}^T_0 = -\frac{1}{4}\,R^2_{\mu\nu} - \frac{1}{4}\,B^2_{\mu\nu} + \frac{M_T^2}{2}\left(R^2_\mu + B^2_\mu\right), \tag{10}$$

where the mass terms take the usual *local* form in this representation.

The new fields automatically obey the Lorentz conditions $\partial^\mu R_\mu = 0$ and $\partial^\mu B_\mu = 0$ and can be considered as usual spin-1 fields. The only difference to the gauge fields consists in the different Lorentz structure of the trilinear interactions

$$\mathcal{L}^T_Y = t\hat{\partial}_\nu\left(\bar{\psi}\sigma^{\mu\nu}\psi\right)R_\mu + it\hat{\partial}_\nu\left(\bar{\psi}\sigma^{\mu\nu}\gamma^5\psi\right)B_\mu \tag{11}$$

and that both R_μ vector and B_μ axial-vector bosons are characterized by the same coupling constant t and the same mass M_T. Therefore, these interactions automatically possess the chiral symmetry under transformations

$$\begin{pmatrix} R_\mu \\ B_\mu \end{pmatrix} \rightarrow \begin{pmatrix} \cos 2\theta & \sin 2\theta \\ -\sin 2\theta & \cos 2\theta \end{pmatrix} \begin{pmatrix} R_\mu \\ B_\mu \end{pmatrix}. \tag{12}$$

On the other hand in order to maintain the chiral symmetry of the Lagrangian (2) under the chiral transformations

$$\begin{pmatrix} S \\ P \end{pmatrix} \rightarrow \begin{pmatrix} \cos 2\theta & \sin 2\theta \\ -\sin 2\theta & \cos 2\theta \end{pmatrix} \begin{pmatrix} S \\ P \end{pmatrix}, \tag{13}$$

one needs to demand an equality of the coupling constants $g_S = g_P$ and the masses $M_S = M_P$ of the spin-0 bosons.

The exchange of massive bosons with a momentum transfer q_μ defines all possible chirally invariant effective four-fermion interactions in the Born approximation

$$\begin{aligned} \mathcal{L}_{\rm eff} &= \frac{g_S^2}{2(M_S^2 - q^2)}\;\bar{\psi}(1+\gamma^5)\psi\;\bar{\psi}(1-\gamma^5)\psi \\ &- \frac{g_V^2}{2(M_V^2 - q^2)}\left(\bar{\psi}\gamma^\mu\psi\right)^2 - \frac{g_A^2}{2(M_A^2 - q^2)}\left(\bar{\psi}\gamma^\mu\gamma^5\psi\right)^2 \\ &- \frac{t^2}{2(M_T^2 - q^2)}\;\bar{\psi}\sigma^{\mu\rho}(1+\gamma^5)\psi\frac{q_\mu q^\nu}{q^2}\bar{\psi}\sigma_{\nu\rho}(1-\gamma^5)\psi. \end{aligned} \tag{14}$$

The specific of the effective tensor interactions consists in the peculiar momentum dependent factor $q_\mu q^\nu/q^2$.

On one hand, this factor ensures good ultraviolet behavior and does not lead to unitarity violation. On the other hand, it has a pole at $q^2 = 0$ and could cause an infrared problem. However, in this paper only the processes with $q^2 \neq 0$ will be considered, I comment shortly its possible solution.

Due to the unusual properties of the chiral fields interactions, the Yukawa coupling constant t as well as the gauge coupling constant g_A exhibits asymptotically free ultraviolet behavior [17], but the region of the low-energy momentum transfers is governed by a nonperturbative physics. One possible solutions to $1/q^2$ pole problem is the Leutwyler's approach [20] to the confinement in presence of constant selfdual abelian background field B. So, nonperturbative effects may modify the propagator in such a way $[1 - \exp(-q^2/B)]/q^2$ to remove singularity at $q^2 = 0$.

It is interesting to compare the angular distributions of the resonance particle-antiparticle s-channel scattering via different intermediate bosons. So the (pseudo)scalar spinless bosons do not prefer any direction in space and lead to the isotropic distribution

$$\frac{{\rm d}\sigma_S}{{\rm d}\cos\vartheta} \propto 1, \tag{15}$$

where ϑ is the scattering angle in the center-of-mass system between incoming and outcoming particles. While

the exchange of the (axial-)vector spin-1 bosons lead to the well known distribution

$$\frac{\mathrm{d}\sigma_V}{\mathrm{d}\cos\vartheta} \propto 1+\cos^2\vartheta \tag{16}$$

up to the linear $\cos\vartheta$ term for P-parity nonconservation.

The tensor interactions also occur through the (axial-) vector spin-1 boson resonances, however they lead to a different from the previous case distribution [21]

$$\frac{\mathrm{d}\sigma_T}{\mathrm{d}\cos\vartheta} \propto \cos^2\vartheta. \tag{17}$$

In this case there exists a characteristic plane, perpendicular to the beam axis, where the emission of the final state pairs is forbidden. Needless to say, that events with a large p_T are the main signature of production of new resonances at colliders. Therefore, the detection of chiral bosons with tensor interactions will be a difficult task, due to the *dip* in the rapidity distribution and the absence of the Jacobian *peak* [10].

Another model independent feature of the tensor interactions is the absence of the interference with the ordinary gauge interactions in the case of massless fermions. In the real case of light fermions negligible interference with the known gauge interactions also makes their detection in low-energy and collider experiments difficult and allows them to escape the experimental constraints.

The bilinear fermion combinations (1) define the quantum numbers J^{PC} for corresponding boson states S, P, V_μ, A_μ, R_μ and B_μ as 0^{++}, 0^{-+}, 1^{--}, 1^{++}, 1^{--} and 1^{+-}, respectively. All these quantum numbers can be assigned to the existing quark-antiquark meson states (see Table I).

J^{PC}	0^{++}	0^{-+}	1^{--}	1^{++}	1^{+-}
$I=0$	f_0	η, η'	ω, ϕ, ω', ϕ'	f_1	h_1
$I=1$	a_0	π	ρ, ρ'	a_1	b_1

TABLE I: The quantum number assignments to the isoscalar $I=0$ and isovector $I=1$ neutral meson states.

So, on one hand the new CP-odd chiral boson B_μ exhibits quantum numbers 1^{+-}, which undoubtedly should be assigned to the existing h_1 and b_1 bosons. On the other hand, the quantum numbers 1^{--} of the new CP-even chiral boson R_μ coincide with the quantum numbers of the vector boson V_μ. Therefore, they could be mixed leading to assignment of the physical states with the quantum numbers 1^{--} in pairs: $\omega-\omega'$, $\phi-\phi'$ and $\rho-\rho'$. However, direct coupling of the *chirally neutral* vector boson V_μ to the *chiral charged* boson R_μ is forbidden by the chiral symmetry and can be realized only as a result of a spontaneous symmetry breaking $\langle S\rangle_0 \neq 0$ through *chirally invariant* trilinear boson interaction

$$\mathcal{L}_\chi = g_\chi\left(S T_{\mu\nu} + P\tilde{T}^{\mu\nu}\right)F^{\mu\nu}, \tag{18}$$

where $F_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu$.

A corresponding model has been developed in [22], where a simple explanation of the dynamic properties of the spin-1 mesons and new mass relations among them have been derived. The results of this approach are in good agreement with the QCD sum rules, the lattice calculations and the experimental data. Let us imagine now, in accordance with the technicolor idea, that analogous phenomenon may be extrapolated to the high energy physics at the Fermi scale, where along with the gauge spin-1 electroweak bosons γ, Z, $W^\pm$ and the spin-0 Higgs bosons H, additional spin-1 *chiral* bosons T are presented.

Since the chiral properties of the new spin-1 bosons are like these of the Higgs bosons, they should come as doublets $T_\mu = \left(T^+_\mu\ T^0_\mu\right)$, taking into account the $SU(2)_L\times U(1)_Y$ symmetry of the SM. However, the chirally invariant trilinear interactions of the type (18) and the gauge interactions (8) lead to new chiral anomalies [18]. In order to cancel the anomalies, additional doublets $U_\mu = \left(U^0_\mu\ U^-_\mu\right)$ with opposite hypercharge to T_μ are introduced. This concerns the Higgs bosons as well, which should also be doubled $H_1 = \left(H^+_1\ H^0_1\right)$ and $H_2 = \left(H^0_2\ H^-_2\right)$.

Therefore, in comparison with the SM, additional boson degrees of freedom are introduced. In spin-0 sector, besides the light SM Higgs boson h, the neutral CP-even H, CP-odd A and the charged $H^\pm$ bosons should be present as in the Minimal Supersymmetric Standard Model (MSSM). The new chiral bosons add eight more spin-1 states: the neutral CP-even $T^R_\mu = \sqrt{2}\,\mathrm{Re}T^0_\mu$, $U^R_\mu = \sqrt{2}\,\mathrm{Re}U^0_\mu$, CP-odd $T^I_\mu = \sqrt{2}\,\mathrm{Im}T^0_\mu$, $U^I_\mu = \sqrt{2}\,\mathrm{Im}U^0_\mu$ and the charged $T^\pm_\mu$, $U^\pm_\mu$ bosons. As far as the spin-0 sector is covered by investigation in the MSSM, I concern only the spin-1 sector of the new chiral bosons.

The trilinear chiral couplings of the new chiral bosons to the quarks $Q^a = (u^a_L\ d^a_L)$, u^a_R, d^a_R and leptons $L^a = (\nu^a_L\ e^a_L)$, ν^a_R, e^a_R

$$\begin{aligned}\mathcal{L}^T_Y = &\left[t^q_{ab}\left(\bar{Q}^a\sigma^{\alpha\beta}d^b_R\right) + t^\ell_{ab}\left(\bar{L}^a\sigma^{\alpha\beta}e^b_R\right)\right]\begin{pmatrix}\hat{\partial}_\alpha T^+_\beta\\ \hat{\partial}_\alpha T^0_\beta\end{pmatrix}\\ &+\left[u^q_{ab}\left(\bar{Q}^a\sigma^{\alpha\beta}u^b_R\right) + u^\ell_{ab}\left(\bar{L}^a\sigma^{\alpha\beta}\nu^b_R\right)\right]\begin{pmatrix}\hat{\partial}_\alpha U^0_\beta\\ \hat{\partial}_\alpha U^-_\beta\end{pmatrix}\\ &+\text{h.c.},\end{aligned} \tag{19}$$

are uniquely fixed by the symmetries. Here a, b are generation indexes and t^q_{ab}, t^ℓ_{ab}, u^q_{ab}, u^ℓ_{ab} are in general arbitrary Yukawa coupling constants. In contrast to ref. [18], additional couplings to the right-handed neutrino states, which are missing from the SM, are included for generality. The latter could be important for the neutrino physics beyond the SM such as leptogenesys and oscillations. However, in the following I suggest that the right-handed neutrino states are very heavy ($m_{\nu_R} \gg m_t$) and have decoupled at the electroweak scale, i.e. one can put $u^\ell_{ab} = 0$.

The explicit form of the interactions (19) gives us the possibility to derive low-energy effective Lagrangian and to investigate the signature of the chiral bosons productions at high energies. These topics with their phenomenological consequences are considered in the next sections. In order to be more definite in the predictions, some model-dependent simplifications are made.

III. LOW-ENERGY INDICATIONS OF NEW CHIRAL INTERACTIONS

In this section I present phenomenological signatures of the chiral bosons effects on the low-energy physics. At the beginning I assume that the chiral bosons are very heavy to be directly produced at low energies, and their effect is only possible through their virtual exchange. For this purpose the effective Lagrangian approach is used. The possibility of the light (even massless) chiral bosons is also viable, but in this case one deals with their unnaturally small Yukawa coupling constants in order to avoid the experimental constraints.

The chiral tensor interactions are topologically similar to the electroweak gauge interactions and, therefore, should contribute to all electroweak processes. However, it is difficult to detect them experimentally on the SM background and some guiding principle is necessary in order to distinguish them from the known interactions. For example, neutral weak currents were detected in the deep-inelastic electron scattering through measurements of P-odd quantities. In the case of the new tensor interactions their *chiral* properties have the main importance.

There are processes, like pion weak decays, where the SM transitions are chirally suppressed and the tensor interactions can manifest themselves at full strength. However, the matrix element of the tensor quark current $\langle 0|\bar{q}\sigma^{\mu\nu}q|\pi\rangle$ is zero by kinematic reasons, and there is no direct contribution to the pion decay from the chiral boson exchange as in the case of the (pseudo)scalar Higgs bosons [23]. This allows the tensor interactions to escape severe experimental constraints. Nevertheless, they can contribute indirectly in the pion decay through the electromagnetic radiative corrections [24] and directly in the radiative pion decay. Namely, the experimental anomalies, which have been observed in the radiative pion decays [25], serve us for the construction of effective new tensor interactions.

Let us consider the charged current transitions, where the only background for the new interactions are the ordinary weak interactions. Searching for deviations from the SM in the neutral current processes on the huge background from the electromagnetic and weak interactions is a more challenging task. However, due to very precise experiments in the determination of chirally suppressed quantities as the anomalous magnetic moments for the electron and muon, it is possible to obtain some evidence for new physics in the neutral sector of the SM, too. This topic will be discussed at the end of this section.

Before proceeding with quantitative calculations it is necessary to fix the arbitrary (in general) Yukawa coupling constants in (19). At present, there is no cogent principle to do this. The simplest but, of course, not unique solution, is to assume quark-lepton and family universality of the tensor interactions

$$t^q_{ab} = t^\ell_{ab} = t\,\delta_{ab}, \qquad u^q_{ab} = u\,\delta_{ab}. \tag{20}$$

This suggestion and the additional hypothesis about a dynamical generation of kinetic terms for the bosons lead to the following relations among the new coupling constants [10]

$$t = \frac{\sqrt{3}}{2}\,u = g, \tag{21}$$

where g is $SU(2)_L$ gauge coupling constant. To convince ourselves in the correctness of this assumption let us note that an analogous relation among Yukawa coupling constants of the various hadron meson resonances in the quark model has successful phenomenological applications [22].

In order to obtain the effective low-energy tensor interactions in the limit of heavy chiral bosons it is necessary to assume the pattern of the chiral symmetry breaking. The most general mass term for the two charged bosons has the form

$$\mathcal{M}^2 = \left(T^+_\alpha \;\; U^+_\alpha \right) \begin{pmatrix} M^2 & -\mu^2 \\ -\mu^2 & m^2 \end{pmatrix} \begin{pmatrix} T^-_\alpha \\ U^-_\alpha \end{pmatrix} \tag{22}$$

with the only requirement of positivity of the determinant $\Delta = M^2 m^2 - \mu^4 = M^2_L M^2_H > 0$ of the square mass matrix, where M^2_L and M^2_H are its eigenvalues corresponding to the lighter and heavier physical mass states. Then the effective tensor quark-lepton interactions read

$$\begin{aligned}\mathcal{L}^{\rm eff}_T =& -\sqrt{2} f_T G_F\, \bar{u}\sigma_{\alpha\rho} d_L\, \frac{4q^\rho q_\beta}{q^2}\, \bar{e}\sigma^{\alpha\beta}\nu_L \\ & -\sqrt{2} f'_T G_F\, \bar{u}\sigma_{\alpha\rho} d_R\, \frac{4q^\rho q_\beta}{q^2}\, \bar{e}\sigma^{\alpha\beta}\nu_L + \text{h.c.},\end{aligned} \tag{23}$$

where

$$f_T = \frac{2M^2_W \mu^2}{\sqrt{3} M^2_L M^2_H} > 0, \qquad f'_T = \frac{M^2_W m^2}{M^2_L M^2_H} > 0 \tag{24}$$

are positive dimensionless coupling constants, which determine the strength of the new tensor interactions relative to the ordinary weak interactions.

The experimental data on the radiative pion decays $\pi \to e\nu\gamma$ [25] show big $\mathcal{O}(10\%)$ deficit of events in the high-E_γ/low-E_e kinematic region. As a matter of fact both terms in (23) with the coupling constants of the order of 10^{-2} could explain the deficit just in the same region [26] through destructive interference of the tensor interactions with the inner bremsstrahlung radiation. However, owing to the electromagnetic radiative corrections, the P-odd tensor quark current $q^\beta \bar{u}\sigma_{\alpha\beta}\gamma^5 d$, which

is present in both terms, leads to a generation of the pseudoscalar quark current $\bar{u}\gamma^5 d$, to which pion decay is very sensitive. Therefore, there is a severe constraint [24] from the pion decay on the coupling constant $|f_T| < 10^{-4}$ or $|f'_T| < 10^{-4}$. In order to avoid this constraint one can assume that the new tensor quark current conserves P-parity due to some unknown symmetry principle and the only P-even tensor quark current $q^\beta \bar{u}\sigma_{\alpha\beta} d$ is present in the effective interactions (23). Because of parity conservation in electromagnetic interactions, it does not contribute to pseudoscalar pion decay. This is realized in the case of equality of the effective coupling constants $f_T = f'_T$ or $\mu^2 = \sqrt{3}m^2/2$.

Taking into account the latter relation, the diagonalization of the mass matrix (22) gives two mass states

$$M^2_{H/L} = \frac{M^2 + m^2 \pm \sqrt{(M^2 - m^2)^2 + 3m^4}}{2}. \qquad (25)$$

If the parameter M^2 is fixed, the maximum value of the lightest mass state is reached at $m^2 = M^2/2$, which defines the physical mass $M^2_L = M^2/4$. It corresponds to an energetically favored exchange by this particle. Then, the heavier state has a mass $M^2_H = M^2 + M^2_L = 5M^2_L$. Accepting the value of the effective tensor coupling $f_T \approx 10^{-2}$, which can explain the deficit of events in the radiative pion decay, one can evaluate the masses of the charged chiral bosons

$$M_H = \sqrt{\frac{2}{f_T}}\, M_W \approx 1137 \text{ GeV}, \quad M_L \approx 509 \text{ GeV}, \quad (26)$$

which will be used in the next section for quantitative estimations of their production cross-sections.

All these relations are model-dependent and follow from the expressions (20,21) for the Yukawa coupling constants. Accepting, that they are of the order of the gauge coupling constants, the relative weakness of the new tensor interactions are explained by the larger chiral boson masses in comparison with the gauge bosons (see (24)). It is interesting to note that the heavier boson mass does not depend on the concrete value of the t/u ratio, while the lighter boson mass is sensitive to it. So, in the symmetric case, $t = u = g$, which has place when $u^\ell_{ab} \neq 0$, $M_L = M_H/\sqrt{6} \approx 464$ GeV. In general the following limit on the mass of the lightest charged boson $M_L < M_H/\sqrt{2} \approx 804$ GeV can be obtained.

The same interactions (23) should inevitably contribute to the neutron decay and in particular it should affect the $\lambda \equiv g_A/g_V$ and V_{ud} determination. However, there is no chiral suppression and the effect of the new interactions is on per mille level. Nevertheless, the experimental accuracy at present is already enough to alarm about the problem.

The most precise λ determination follows from the electron asymmetry parameter A measured in the polarized neutron decay. The latest preliminary result of the PERKEO collaboration $A = -0.1195 \pm 0.0004$ [27] leads to a very high absolute value of $\lambda = -1.2755 \pm 0.0011$. Using the PDG value for the neutron lifetime $\tau_n = 885.7 \pm 0.8$ s [28], one can evaluate $V^n_{ud} = 0.97081 \pm 0.00088$. This value is $3.2\,\sigma$ lower than the extracted one from superallowed Fermi nuclear decays $V^F_{ud} = 0.97377 \pm 0.00027$ [29] and may be explained through the presence of the new interactions (23).

Their effect, considered in [30], leads to a negative contribution into A over the whole electron spectrum and to a wrong λ evaluation from experimental data. In the same time in a first approximation they do not distort the recoil proton spectrum and do not contribute to the neutron lifetime. Therefore, it is interesting to compare thus obtained $\lambda(A)$ with $\lambda(a)$ determined from the correlation coefficient a, relying on the method based on measurements of the proton kinetic energy spectrum in the unpolarized neutron decay. I hope the results of the new aSPECT experiment can clarify the problem.

Assuming lepton universality of the tensor interactions (23) one can investigate their effects in the muon and tau lepton decays, as well. They can affect, for example, hadronic τ decays with the same effective coupling constant f_T. By kinematic reason, they do not contribute to the single pion decay channel $\tau \to \pi\nu$. At the same time, the two pion decay channel $\tau \to \rho\nu \to 2\pi\nu$ through the vector ρ meson should be considerably affected on the level of $\mathcal{O}(10\%)$, due to their interference with ρ meson tensor quark current and very big mass of the τ lepton [31]. Using the CVC hypothesis [32] the branching ratio for the two pion decay channel can be predicted from electromagnetic process $e^+e^- \to \gamma^* \to \rho \to 2\pi$, where the tensor interactions (23) are not operative. Indeed, the predicted value at present is $4.5\,\sigma$ lower than the measured two pion branching ratio of the τ decay [33].

The fact that the tensor interactions (23) with the same effective coupling constant f_T can explain simultaneously the destructive interference in the radiative pion decay, the anomalously big asymmetry parameter A in the polarized neutron decay, the excess of two pion production in τ decay and at the same time still can escape the other experimental constraints is a witness for their vitality. The discrepancy between the two pion production in the e^+e^- annihilation and the τ decay causes another problem connected with the prediction of the anomalous muon magnetic moment.

At present, due to the mentioned discrepancy, the tau data is not used for a prediction of the hadron contribution to the anomalous muon magnetic moment, while the e^+e^- data lead to a value [33] by $3.3\,\sigma$ lower than measured one [34]

$$\delta a_\mu = a^{\rm exp}_\mu - a^{\rm th}_\mu = (27.5 \pm 8.4) \times 10^{-10}. \qquad (27)$$

It is a huge discrepancy as far as the contribution of the massive SM weak bosons is only $a^{\rm weak}_\mu = (15.4 \pm 0.2) \times 10^{-10}$. Therefore, naively one can expect a much lower contribution from heavier new bosons, which cannot explain the difference (27).

Could the new neutral chiral bosons T^0_μ or U^0_μ, which couple anomalously to matter (19), explain the difference? Of course, at present it is impossible, having in mind the completely unknown neutral sector of the tensor interactions, to give quantitative value of the effect. However, it is still possible to predict the sign of the difference. The idea is the following.

The electromagnetic interactions of the photon A_α with the charged leptons ℓ read

$$\mathcal{L}^\gamma_{\rm int} = e\,\bar\ell\gamma^\alpha\ell\;A_\alpha + a_\ell\frac{e}{2m_\ell}\;\bar\ell\sigma^{\alpha\beta}\ell\;\partial_\alpha A_\beta, \tag{28}$$

where the first term is the gauge interaction and the second one is the effective photon coupling to the anomalous magnetic moment of the lepton. The chiral structure of the effective anomalous photon coupling coincides with the chiral structure of the new tensor interactions (19)

$$\begin{aligned}\mathcal{L}^N_Y &= \frac{t}{\sqrt2}\left(\bar d\sigma^{\alpha\beta}d + \bar\ell\sigma^{\alpha\beta}\ell\right)\hat\partial_\alpha T^R_\beta \\ &+ i\frac{t}{\sqrt2}\left(\bar d\sigma^{\alpha\beta}\gamma^5 d + \bar\ell\sigma^{\alpha\beta}\gamma^5\ell\right)\hat\partial_\alpha T^I_\beta \\ &+ \frac{u}{\sqrt2}\left(\bar u\sigma^{\alpha\beta}u\right)\hat\partial_\alpha U^R_\beta, \\ &+ i\frac{u}{\sqrt2}\left(\bar u\sigma^{\alpha\beta}\gamma^5 u\right)\hat\partial_\alpha U^I_\beta,\end{aligned} \tag{29}$$

where only the T_α chiral bosons couple to the *charged* leptons.

An additional contribution to the anomalous magnetic moment of the lepton could arise from a mixing between the photon A_α and the chiral boson T^R_α. In the chiral invariant limit there are not mixing between them. However, it can appear as a result of spontaneous symmetry breaking. Simple calculations of the mixing, for example, through the loops with the massive fermions lead to the following contribution

$$\mathcal{L}^{\gamma T}_{\rm mix} = -\kappa^{\gamma T} m^*_d\left(\partial^\alpha A^\beta - \partial^\beta A^\alpha\right)\hat\partial_\alpha T^R_\beta, \tag{30}$$

where

$$\kappa^{\gamma T} = \frac{\sqrt2\,e\,t}{4\pi^2}\sum_{i=\ell,d}\frac{m_i}{m^*_d}\left(\ln\frac{\Lambda^2}{m_i^2} - 1\right), \tag{31}$$

m^*_d is the effective mass of the *down* fermions and Λ is the effective ultraviolet cutoff.

In contrast to the case of the charged chiral bosons, the neutral states T^0_α and U^0_α do not mix between each other, because they couple to the different types, *up* and *down* fermions (29). The CP-odd states T^I_α and U^I_α also decouple from the CP-even states T^R_α and U^R_α in the case of the CP symmetry, due to their quantum numbers.

Therefore, the full free Lagrangian for the neutral bosons reads

$$\begin{aligned}\mathcal{L}_0 &= \frac12\begin{pmatrix} A_\alpha & Z_\alpha & U^R_\alpha & T^R_\alpha\end{pmatrix}\mathcal{K}^{\alpha\beta}_{\rm even}\begin{pmatrix}A_\beta\\ Z_\beta\\ U^R_\beta\\ T^R_\beta\end{pmatrix} \\ &+ \frac12\begin{pmatrix} U^I_\alpha & T^I_\alpha\end{pmatrix}\mathcal{K}^{\alpha\beta}_{\rm odd}\begin{pmatrix}U^I_\beta\\ T^I_\beta\end{pmatrix},\end{aligned} \tag{32}$$

where

$$\mathcal{K}^{\alpha\beta}_{\rm even} = \left(\frac{q^\alpha q^\beta}{q^2} - g^{\alpha\beta}\right)\times \\ \times\begin{pmatrix} q^2 & \kappa^{\gamma Z}q^2 & \kappa^{\gamma U}m^*_u|q| & \kappa^{\gamma T}m^*_d|q| \\ \kappa^{\gamma Z}q^2 & q^2 - M_Z^2 & \kappa^{ZU}m^*_u|q| & \kappa^{ZT}m^*_d|q| \\ \kappa^{\gamma U}m^*_u|q| & \kappa^{ZU}m^*_u|q| & q^2 - M^2_{U_R} & 0 \\ \kappa^{\gamma T}m^*_d|q| & \kappa^{ZT}m^*_d|q| & 0 & q^2 - M^2_{T_R}\end{pmatrix} \tag{33}$$

and

$$\mathcal{K}^{\alpha\beta}_{\rm odd} = \left(\frac{q^\alpha q^\beta}{q^2} - g^{\alpha\beta}\right)\times \\ \times\begin{pmatrix} q^2 - M^2_{U_I} & 0 \\ 0 & q^2 - M^2_{T_I}\end{pmatrix}. \tag{34}$$

The mixings between the gauge bosons and the CP-even chiral bosons (30) are tiny since they are proportional to the light masses of the ordinary fermions and the small coupling constants (31). Therefore, the matrix (33) is almost diagonal and the physical state of the CP-even chiral boson $\mathcal{T}^R_\alpha$, for example, contains small but not negligible admixture of the gauge bosons

$$\mathcal{T}^R_\alpha \simeq T^R_\alpha - \kappa^{\gamma T}\frac{m^*_d|q|}{M^2_{T_R}}A_\alpha - \kappa^{ZT}\frac{m^*_d|q|}{M^2_{T_R} - M^2_Z}Z_\alpha. \tag{35}$$

This leads to an additional positive contribution to the anomalous magnetic moment of the lepton

$$\delta a_\ell \simeq \frac{\sqrt2\,t}{e}\,\kappa^{\gamma T}\frac{m_\ell\,m^*_d}{M^2_{T_R}} > 0. \tag{36}$$

The sign of the contribution is in agreement with the discrepancy in the experimental data for the muon (27).

Assuming the universality of the Yukawa tensor interactions and taking into account the magnitude of the discrepancy (27) it is possible to predict the additional contribution into the anomalous electron magnetic moment

$$\delta a_e \simeq (13.3 \pm 4.1)\times 10^{-12} \tag{37}$$

from the photon mixing with the chiral boson. It is interesting to note that this contribution is well above the non-QED contributions $a_e^{\rm HAD} = 1.671(19)\times 10^{-12}$, $a_e^{\rm EW} = 0.030(01)\times 10^{-12}$ [35] and the experimental error $\delta a_e^{\rm exp} = 0.76\times 10^{-12}$ [36]. Therefore, if it is real, it should give essential correction

$$\delta\alpha^{-1} \simeq (15.7 \pm 4.8)\times 10^{-7} \tag{38}$$

to the determination of the fine structure constant from the anomalous electron magnetic moment using the QED calculations [37]. Unfortunately, independent α determinations [38, 39] have errors comparable with the contribution (38). Therefore, new α measurements are badly needed. This will allow to pin down the problem with the anomalous muon magnetic moment and the presence of the new tensor interactions.

The physical boson masses are not affected so much by the chiral symmetry breaking. Moreover, the photon remains massless protected by the gauge-invariant mixing (30). The masses of the neutral chiral bosons

$$\begin{aligned} M_{T_R} &\simeq M_{T_I} \simeq M \approx 1017 \text{ GeV}, \\ M_{U_R} &\simeq M_{U_I} \simeq m \approx 719 \text{ GeV}, \end{aligned} \tag{39}$$

are expressed through diagonal elements of the mixing mass matrix for the corresponding doublets of the charged chiral bosons (22) up to a negligibly small correction of κ^2. This situation is completely different from the case of the low-lying hadron meson vector states, where the mixing is maximal and the physical masses differ considerably with respect to the chirally symmetric case [19].

IV. CHIRAL BOSON PRODUCTION AT THE FERMILAB TEVATRON

Up to now I discussed manifestations of the tensor interactions in low-energy experiments as hints for the existence of the fundamental intermediate chiral bosons. However, the crucial confirmation for their existence should come from their direct production at the colliders with a unique signature.

From the previous considerations it follows that the mass of the lightest chiral bosons is around 500 GeV. Since they are charged bosons, they could be produced at the lepton colliders only in pairs or in association with other charged boson, like W. The lightest neutral chiral bosons do not interact with leptons and cannot be produced at the lepton colliders at all. Therefore, to produce the pairs of the lightest charged chiral bosons or the heaviest neutral chiral bosons, which mass is just two times bigger than the mass of the lightest bosons, one needs a lepton collider with energy above 1 TeV. The International Linear Collider (ILC) with such energy would be an ideal place to produce and to study these particles.

In general the low-energy effective tensor interactions can be tested at the lepton colliders as the LEP and the SLC or at the Hadron-Electron Ring Accelerator (HERA). Unfortunately, these interactions do not interfere with the SM $(V-A)$ interactions of the light particles and their contribution into the cross-section is of the order of $f_T^2 \sim 10^{-4}$. This is an order of magnitude smaller even than the experimental errors 0.1% at the high-precision lepton colliders. Nevertheless, they can still affect the observables connected with the heavy τ lepton and b quark.

At present only the Tevatron hadron collider at Fermilab is powerful enough to produce and detect with non-negligible probability at least the lightest charged chiral bosons [10]. Let us consider this possibility in more detail. Since the case of the hypothetical W' gauge boson with the SM couplings is well known, in the following I shall compare its properties with the new chiral boson's ones.

Due to the mass mixing (22) the physical states are represented by two orthogonal combinations $\mathcal{U}^\pm_m = (\sqrt{3}\,U^\pm_m + T^\pm_m)/2$ and $\mathcal{T}^\pm_m = (\sqrt{3}\,T^\pm_m - U^\pm_m)/2$, which correspond to light and heavy massive particles, respectively. Their Yukawa interactions take the form

$$\begin{aligned} \mathcal{L}^C_Y = \frac{g}{2}\left(\bar{u}_a\sigma^{mn}d_{Ra} + \bar{\nu}_a\sigma^{mn}e_{Ra}\right)\left(\hat{\partial}_m\mathcal{U}^+_n + \sqrt{3}\,\hat{\partial}_m\mathcal{T}^+_n\right) \\ + g\left(\bar{u}_a\sigma^{mn}d_{La}\right)\left(\hat{\partial}_m\mathcal{U}^+_n - \frac{1}{\sqrt{3}}\,\hat{\partial}_m\mathcal{T}^+_n\right) + \text{h.c.} \end{aligned} \tag{40}$$

Special relations (20,21) between the Yukawa coupling constants assure the same total decay widths into fermions

$$\Gamma^V_{tot} = \frac{g^2}{4\pi} M_V \tag{41}$$

both for the gauge and chiral bosons with the same mass $M_V \gg m_t$. It follows from the dynamical generation of kinetic terms for the bosons using the condition of equality of all one-loop fermion contributions into self-energy bosons, the imaginary part of which is proportional to the decay width for the corresponding bosons. In the following, as a first approximation, I take into account only the fermion decay channels and do not consider boson decays into the known gauge and Higgs bosons.

The relations between the lepton and quark decay widths depend on the boson type and the model

$$\begin{aligned} &\Gamma^{W'}_{ud} = \Gamma^{W'}_{cs} = \Gamma^{W'}_{tb} = 3\,\Gamma^{W'}_{e\nu_e} = 3\,\Gamma^{W'}_{\mu\nu_\mu} = 3\,\Gamma^{W'}_{\tau\nu_\tau}, \\ &\Gamma^{\mathcal{U}}_{ud} = \Gamma^{\mathcal{U}}_{cs} = \Gamma^{\mathcal{U}}_{tb} = 15\,\Gamma^{\mathcal{U}}_{e\nu_e} = 15\,\Gamma^{\mathcal{U}}_{\mu\nu_\mu} = 15\,\Gamma^{\mathcal{U}}_{\tau\nu_\tau}, \\ &\Gamma^{\mathcal{T}}_{ud} = \Gamma^{\mathcal{T}}_{cs} = \Gamma^{\mathcal{T}}_{tb} = \frac{13}{3}\,\Gamma^{\mathcal{T}}_{e\nu_e} = \frac{13}{3}\,\Gamma^{\mathcal{T}}_{\mu\nu_\mu} = \frac{13}{3}\,\Gamma^{\mathcal{T}}_{\tau\nu_\tau}, \end{aligned} \tag{42}$$

or

$$\begin{aligned} \Gamma^{W'}_{tot} &= 4\,\Gamma^{W'}_{ud} = 12\,\Gamma^{W'}_{\ell\nu}, \\ \Gamma^{\mathcal{U}}_{tot} &= \frac{16}{5}\,\Gamma^{\mathcal{U}}_{ud} = 48\,\Gamma^{\mathcal{U}}_{\ell\nu}, \\ \Gamma^{\mathcal{T}}_{tot} &= \frac{48}{13}\,\Gamma^{\mathcal{T}}_{ud} = 16\,\Gamma^{\mathcal{T}}_{\ell\nu}. \end{aligned} \tag{43}$$

Here I accept a diagonal CKM mixing matrices for the W' and the chiral bosons.

The resonant production cross-section of the intermediated bosons at parton level is proportional to their partial decay width into the quarks pair

$$\begin{aligned} \hat{\sigma}(u\bar{d} \to V^+) &= \frac{4\pi^2}{3M_V}\,\Gamma^V_{ud}\,\delta(\hat{s} - M_V^2) \\ &= \frac{\pi g^2}{3}\mathcal{B}(V^+ \to u\bar{d})\,\delta(\hat{s} - M_V^2), \end{aligned} \tag{44}$$

where $\hat{s} = (p_u + p_{\bar{d}})^2$ is the invariant Mandelstam variable. Here and in the following I shall denote with a hat the variables in the parton center-of-mass system.

Neglecting the small contribution of the see quarks at the Tevatron $p\bar{p}$ collider, the total production cross-section reads

$$\sigma_V^{\rm Tev} = \frac{\pi g^2}{3s}\mathcal{B}(V \to u\bar{d})\int_\tau^1 u(x, M_L) d\left(\frac{\tau}{x}, M_L\right)\frac{{\rm d}x}{x}, \quad (45)$$

where s is the square of the center-of-mass energy and $\tau = M_V^2/s$. From the relations (43) and (41) it follows that the quark branching ratios $\mathcal{B}(W' \to u\bar{d}) \simeq 25.0\%$, $\mathcal{B}(\mathcal{U} \to u\bar{d}) \simeq 31.3\%$ and $\mathcal{B}(\mathcal{T} \to u\bar{d}) \simeq 27.1\%$ do not differ very much for the different type of bosons, therefore, the total production cross-sections of the gauge and the chiral bosons with the same mass have approximately the same magnitudes.

However, the cleanest way to detect the production of the heavy intermediate bosons at the hadron colliders is to look for their practically backgroundless decay channels into leptons. It is interesting to note that the lepton branching of the lightest chiral boson $\mathcal{B}(\mathcal{U} \to \ell\bar{\nu}) \simeq 2.1\%$ is considerably smaller than the others $\mathcal{B}(W' \to \ell\bar{\nu}) \simeq 8.3\%$ and $\mathcal{B}(\mathcal{T} \to \ell\bar{\nu}) \simeq 6.3\%$. Therefore, its leptophobic character leads to considerably small cross-section in the lepton channel (see Fig.1).

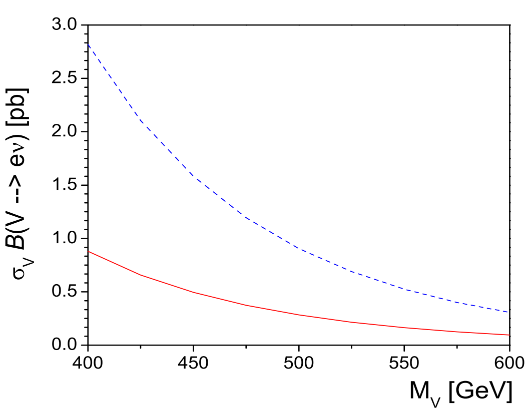

FIG. 1: The production cross-sections of the gauge W' boson (dashed) and the chiral $\mathcal{U}$ boson (solid) as functions of their masses.

For these calculations the parton distribution functions CTEQ6M [40] and the factor $K = 1$ have been used. So, for the reference mass $M_V = M_L \approx 509$ GeV (26) of the lightest chiral boson, the corresponding cross sections are

$$\begin{aligned} \sigma_{W'} \times \mathcal{B}(W' \to \ell\bar{\nu}) &\approx 0.82 \text{ pb}, \\ \sigma_{\mathcal{U}} \times \mathcal{B}(\mathcal{U} \to \ell\bar{\nu}) &\approx 0.26 \text{ pb}. \end{aligned} \quad (46)$$

The latter cross-section is by factor 16/5 smaller than the former one. This fact makes the detection of the chiral bosons in the lepton channels more difficult than the gauge ones.

But not only this fact prevents the discovery of the new chiral bosons up to now. Another unusual and unexpected feature of the chiral bosons, connected to their anomalous interactions (19) with fermions, has place. Let us compare the normalized angular distributions of the leptons from the decays of the gauge W^-

$$\frac{{\rm d}N_W}{{\rm d}\Omega} = \begin{cases} \frac{3}{16\pi}(1 \mp \cos\theta^*)^2, & \lambda = \pm 1, \\ \frac{3}{8\pi}\sin^2\theta^*, & \lambda = 0, \end{cases} \quad (47)$$

and the chiral $\mathcal{U}^\pm$

$$\frac{{\rm d}N_{\mathcal{U}}}{{\rm d}\Omega} = \begin{cases} \frac{3}{8\pi}\sin^2\theta^*, & \lambda = \pm 1, \\ \frac{3}{4\pi}\cos^2\theta^*, & \lambda = 0, \end{cases} \quad (48)$$

bosons, where θ^* is the angle of the lepton with respect to the longitudinal axis in the boson-rest frame and λ is the boson helicity.

For example, the left-handed quark d (from the proton) interacting with the right-handed anti-quark $\bar{u}$ (from the anti-proton) can produce W^- with spin projection on the proton beam direction -1. Hence, the decay leptons are distributed as $(1 + \cos\theta^*)^2$. While chiral particle production arises from the interaction of a quark and an anti-quark with the same helicities, that leads to zero helicity of the produced chiral boson and the backward-forward symmetric $\cos^2\theta^*$ lepton distribution [21].

Indeed, the cross-section for $p + \bar{p} \to \mathcal{U} + X \to \ell + X'$ process

$$d\sigma = \frac{1}{3}\int {\rm d}x_1 {\rm d}x_2\, u(x_1, M_L) d(x_2, M_L)\, {\rm d}\hat{\sigma}(\hat{s}, \hat{t}) \quad (49)$$

is expressed through the relevant differential cross-section of the parton subprocess $d + \bar{u} \to \mathcal{U}^- \to \ell + \bar{\nu}$

$$E_\ell \frac{{\rm d}^3\hat{\sigma}(\hat{s}, \hat{t})}{{\rm d}^3 p_\ell} = \frac{5g^4}{(32\pi)^2}\frac{(\hat{s} + 2\hat{t})^2\delta(\hat{s} + \hat{t} + \hat{u})}{\hat{s}|\hat{s} - M_L^2 + iM_L\Gamma_{tot}^{\mathcal{U}}|^2}, \quad (50)$$

where $\hat{s} = (p_d + p_{\bar{u}})^2$, $\hat{t} = (p_{\bar{u}} - p_\ell)^2$ and $\hat{u} = (p_d - p_\ell)^2$ are the Mandelstam variables. In the center-of-mass parton system the differential cross-section shows the following distribution

$$\frac{{\rm d}\hat{\sigma}}{{\rm d}\Omega} \propto (\hat{s} + 2\hat{t})^2 \propto \cos^2\hat{\theta} = 1 - \frac{4\hat{p}_T^2}{\hat{s}}. \quad (51)$$

Here $\hat{\theta}$ is the angle between the lepton and the parton direction, which coincide with the angle θ^* in the boson-rest frame, and $\hat{p}_T^2$ is the square of the transverse lepton momentum.

Since the latter is invariant under longitudinal boosts along the beam direction, the distribution (51) versus $\hat{p}_T$ holds also in the lab frame $p_T = \hat{p}_T$. Changing variables in the differential cross-section from $\cos\hat{\theta}$ to p_T^2

$$\frac{{\rm d}\cos\hat{\theta}}{{\rm d}p_T^2} = -\frac{2}{\hat{s}}\left(\sqrt{1 - \frac{4p_T^2}{\hat{s}}}\right)^{-1} \quad (52)$$

leads to a kinematical singularity at the endpoint $p_T^2 = \hat{s}/4$, which gives the prominent Jacobian peak in the W decay distribution.

In contrast to this, the pole in the decay distribution of the chiral bosons is cancelled out and, moreover, the distribution reaches zero at the endpoint $p_T^2 = \hat{s}/4$. The chiral boson decay distribution has a broad smooth bump with a maximum below the kinematical endpoint, instead of a sharp Jacobian peak (Fig. 2). In the case when the chiral boson is produced with no transverse momentum, the transverse mass of the lepton pair is related to p_T as $M_T(\ell\bar{\nu}) = 2p_T$ and the Jacobian peak is absent in M_T distribution as well.

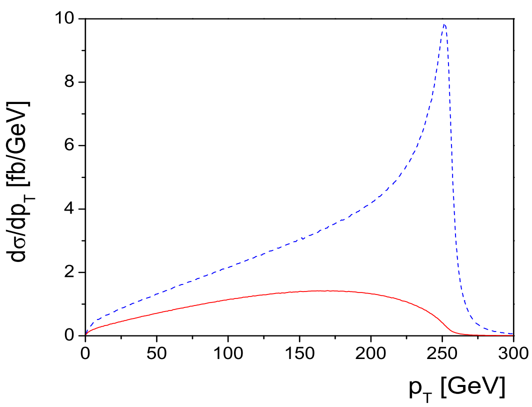

FIG. 2: The differential cross-sections for the gauge W' boson (dashed) and the chiral $\mathcal{U}$ boson (solid) as functions of the lepton transfers momentum.

Therefore, the transverse momentum/mass decay distribution of the chiral bosons differs drastically from the distribution of the gauge bosons. Even relatively small decay width of the chiral bosons leads to a wide distribution, that obscures their identification as resonances at hadron colliders. The form of the decay distribution for the chiral bosons resembles the bump anomalies in the inclusive jet E_T distribution, reported by the CDF Collaboration [41] many years ago. Although this problem has been solved in the framework of the SM by changing the gluon distribution functions [42], it could be reconsidered in the light of the new form of the decay distribution as a real physical signal from decays of different chiral bosons, both charged and neutral.

Analysing the bumps in the jet transverse energy distribution in Fig. 1 of ref. [41], we can find the endpoint of the first bump at 250 GeV and guess about the second bump endpoint from the minimum around 350 GeV. If we assign the first bump to the hadron decay products of the lightest charged bosons, which exactly corresponds to the estimated mass from eq. (26), the second endpoint hints to a mass around 700 GeV, which is also in a quantitative agreement with our estimations (39) for the mass of the lightest leptophobic neutral boson. However, taking into account the large systematic uncertainties in jet production, these conclusions may be premature, unless they are confirmed in other channels.

In the following the CalcHEP [13] package will be used for the numeric calculations of various distributions. For these purposes I have introduced the new chiral bosons in the package and implemented the corresponding interactions for them. Indeed, the CalcHEP model does not allow to introduce the interactions (19) directly. The problem is connected with unusual $1/\sqrt{q^2}$ momentum dependence of the Yukawa interactions, when the chiral bosons are described by the four-dimensional Lorentz vector V_α with the usual propagator $-ig_{\alpha\beta}/(q^2 - M_V^2)$. It can be avoided for the processes with the chiral particles on the mass shell $q^2 = M_V^2$, when the the function $1/\sqrt{q^2}$ can be replaced by constant factor $1/M_V$, but not for the intermediate chiral bosons.

The solution of the problem was prompted to me by A. Pukhov, one of the authors of the CalcHEP package. It consists in introducing of pair of massive particle and its massless ghost in such a way that the propagator for the intermediate states multiplied by the momentum dependent factor $1/q^2$, which comes from the Yukawa couplings, can be represented as a difference of two propagators

$$\frac{1}{q^2}\,\frac{1}{q^2 - M_V^2} = \frac{1}{M_V^2}\left(\frac{1}{q^2 - M_V^2} - \frac{1}{q^2}\right) \qquad (53)$$

with the constant factor $1/M_V^2$. Therefore, the situation is reduced to the previous case of the chiral boson description on the mass shell by effective Yukawa interactions with the constant factor $1/M_V$. This dimension factor naively indicates a problem with high-energy behaviour of scattering amplitudes and renormalizability of the model, which are restored by the including of the ghosts.

The distributions in the Fig. 2 have been calculated without any kinematical cuts. However, experimental detectors always have dead zones where the particles cannot be registered. The simple examples are the backward-forward regions along the beam direction at the colliders. The CaclHEP package with its effective Monte Carlo integration allows simply and quickly to calculate divers distributions with experimental cuts as a first approximations to the detector simulations.

The cuts in the backward-forward regions lead to miss in a essential part of the events from the chiral boson decays due to the mentioned previously specific angular distribution. So, in the Fig. 3 the differential cross-sections of the gauge W' boson and the chiral $\mathcal{U}$ boson as functions of the lepton pseudorapidity are shown. While the maximum of the gauge boson distribution is centered at the small lepton pseudorapidities, which correspond to the central part of the detector, the chiral boson distribution has minimum in this region and its maxima are placed at the edges of the CDF and D0 central calorimeters. Based on the fact that the major part of the leptons stemming

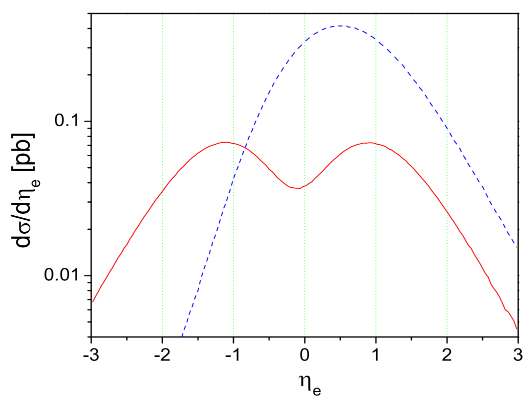

FIG. 3: The differential cross-sections for the gauge W' boson (dashed) and the chiral $\mathcal{U}$ boson (solid) as functions of the lepton pseudorapidity.

from the W' decays are emitted in the central detector region, both collaborations have analyzed the spectrum of the transverse high-energy electrons only in the central electromagnetic calorimeters $|\eta_e| \leq \eta_{cut} \simeq 1$. The ratio $R = (N_{tot} - N_{mis})/N_{tot}$, where N_{mis} is the number of missing due to the cut events, for the W' boson and the chiral $\mathcal{U}$ boson is shown in the Fig. 4.

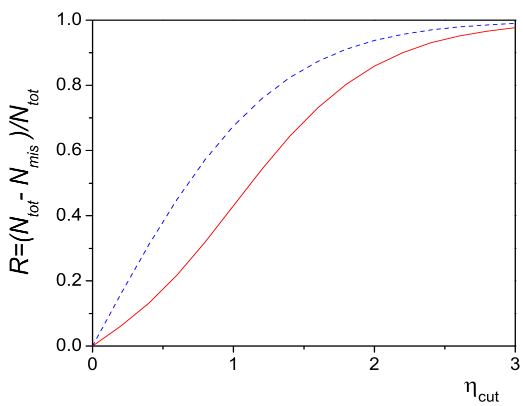

FIG. 4: The ratio of the detected with $|\eta_e| \leq \eta_{cut}$ events to their total number for the gauge W' boson (dashed) and the chiral $\mathcal{U}$ boson (solid).

As seen from Fig. 4 the curve for the chiral $\mathcal{U}$ boson lies always under the W' curve, and at $\eta_{cut} \simeq 1$ there are around 70% detected events in the case of the gauge W' boson and only 45% in the case of the chiral $\mathcal{U}$ boson.

If the new chiral boson production takes place with subsequent decay into the lepton and its antineutrino, an eventual excess should be watched in the region 350 GeV $< M_T <$ 500 GeV, where the background from the tail of the W decays is considerably small. Indeed, such an excess about 2σ has been pointed out recently by the CDF Collaboraion [43] in the same region. In the case of the gauge W' boson production an excess should be peaked around $M_T \approx 500$ GeV or $p_T \approx 250$ GeV. The latter was unambiguously rejected by the D0 data [44] with more better calorimetry than the CDF detector (Fig. 5). At

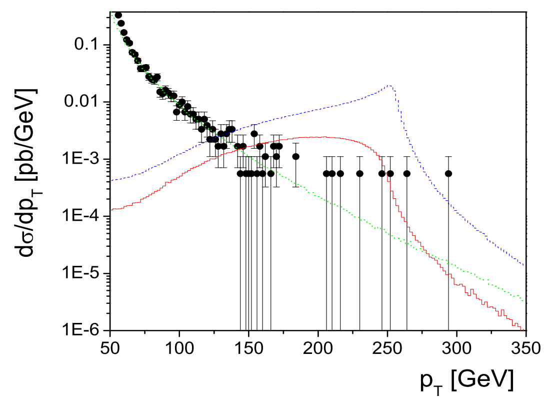

FIG. 5: The differential cross-sections for the gauge W (dotted), W' (dashed) and the chiral $\mathcal{U}$ (solid) bosons as a function of the electron transverse momentum versus the D0 data.

the same time the chiral boson distribution, probably up to some common normalized factor, is still in agreement with the D0 experimental data. The excess should look like as a shoulder rather than a peak and the small number of the events in this region cannot give conclusive statement about the excess.

Nevertheless, the $p\bar{p}$ colliders such as the CERN SPS and the Fermilab Tevatron have unique possibility to measure lepton asymmetry, which cannot be done at the powerful but pp symmetric LHC machine. So, one of the crucial confirmations of the W properties at the SPS, even with a small number of events, was revealing the backward–forward asymmetry in the angular distribution of the lepton emission angle θ^* in the W-rest frame [45]. Application of analogous analysis for the events from the W tail at the Tevatron would allow to reduce additionally the background in the search of the gauge right-handed W'_R and the chiral $\mathcal{U}$ bosons.

Such analysis is even more important in the search of heavy boson production through their hadronic decay channels, where the background from the strong interactions is overwhelming. So, the cross-sections for the $t\bar{b}$ quark channel are

$$\begin{aligned} \sigma_{W'} \times \mathcal{B}(W' \to t\bar{b}) &\approx 1.91 \text{ pb}, \\ \sigma_{\mathcal{U}} \times \mathcal{B}(\mathcal{U} \to t\bar{b}) &\approx 3.56 \text{ pb}, \end{aligned} \tag{54}$$

for the intermediate states with the W' and $\mathcal{U}$ bosons, correspondingly. The latter cross-section is 25/16 times bigger than the former one, due to the hadrophilic character of the chiral boson. Therefore, its detection in the

hadronic channels could be even easy than the detection of the new gauge boson.

While the light quark decay channels are swamped by multijet background, the $t\bar{b}$ pair of the heavy b quark and the short living t quark with its subsequent decay to Wb pair allow to make jet b-tagging, where one of the jets must have a displaced secondary vertex. Searching for the intermediate heavy bosons in this channel has been fulfilled by both the D0 and CDF collaborations, as for this purpose the part of the same datasets of the single top production analyses has been used. Owing to their high masses this analysis is even simpler than the single top production searches, because at such energies the background is considerably reduced.

Indeed, the CDF collaboration even with approximately 1 fb^{-1} of the $p\bar{p}$ RUN II data is still unable to pin down the single top quark production, while a slight excess in the region near $450-500$ GeV in the invariant mass of the reconstructed W and two leading jets (M_{Wjj}) has been pointed out in [46]. The excess is seen in the 2-jet mass distribution histogram with 40 GeV bin's width. It is also expected to find a Jacobian peak in the transverse momentum distribution of the leading b jet or reconstructed top, however, if this excess is due to the chiral boson production, it could not be there.

The common wisdom, that a peak in the invariant mass distribution of the two final particles must correspond to the peaks in their transverse mass distributions M_T at the same value, is not valid for the chiral bosons. On the other hand the peaks in the invariant mass distribution come from the Breit–Wigner propagator form, which is the same as for the gauge and chiral bosons in the Born approximation. However, some small different corrections should be applied to the shape of the resonance curve for the chiral bosons, due to their different couplings to the fermions.

It is interesting to note that this excess is in some sense a confirmation of the excess in the leptonic channel [43] of the same collaboration. Therefore, an independent result from the D0 collaboration is very important. Their published result [47] is based on 230 pb^{-1} of integrated luminosity and does not show any excess in the histogram with the bin's width of 50 GeV. However, it should always be taken into account that the narrow peak could be missed due to the smearing effect of the detector resolution or an insufficient statistic. Indeed, the right histogram in the Fig. 3 of the conference paper [48] of the same collaboration with the bin's width of the 45 GeV reveals, nevertheless, the weak peak in the same region of the 500 GeV. All these not statistically significant results for the separated analyses may give a more conclusive answer after their combining and additional investigation of the angular distributions of the events in this region. The only difficulty could arise from reconstructing the full kinematics due to double solution for the neutrino longitudinal momentum.

Concluding this section I would like to note that some hints for the existence of the lightest charged chiral boson already exist in the Tevatron data. What about signature of the other chiral bosons? The next to the lightest chiral boson are the two completely leptophobic neutral CP-even U^R and CP-odd U^I bosons with approximately the same mass $m \approx 719$ GeV, which couple only to *up*-type quarks. Therefore, they could be looked for in the $t\bar{t}$ decay channel. However, the small cross-section

$$\sigma_U \times \mathcal{B}\,(U \to t\bar{t}) \approx 1.63 \text{ pb} \tag{55}$$

still hinders their revealing.

The other pair of more heavy neutral chiral bosons, CP-even T^R and CP-odd T^I can be seen in their dilepton decay channels. However, their high masses $M \approx 1017$ GeV lead to completely negligible cross-section

$$\sigma_T \times \mathcal{B}\left(T \to \ell\bar{\ell}\right) \approx 2.3 \text{ fb} \tag{56}$$

at the Tevatron. The heaviest charged chiral boson $\mathcal{T}$ with mass $M_H \approx 1137$ GeV also leads to negligible cross-section in the lepton channel

$$\sigma_{\mathcal{T}} \times \mathcal{B}\,(\mathcal{T} \to \ell\bar{\nu}) \approx 2.2 \text{ fb} \tag{57}$$

at the Tevatron.

V. CERN LHC PROSPECTS AND CONCLUSIONS

The LHC at CERN belongs to the next generation of hadron colliders. There is no doubt that it will be discovery machine with its 14 TeV collider energy in the pp center-of-mass. All new particles with non-negligible coupling constants to the ordinary matter up to a mass of the order of $2-3$ TeV could be discovered if they exist. The new chiral bosons, with coupling constants of the order of the gauge ones and the predicted masses, comfortably fall into this category. Nevertheless, the unusual properties of the chiral bosons, which are still unknown to the experimentalists, require more work in simulations.

As a first step into this direction the corresponding model was introduced in the CalcHEP package. In the agreement with the Les Houches Accord [49] the package has an interface with the PYTHIA. However, the experimental software on simulation and reconstruction (see, for example, [50]) often has itself an interface with the event generator programs like PYTHIA. Therefore, I have built the code for the chiral bosons directly into the PYTHIA, as well. However, the right way to do this is to use user-defined external processes machinery. In order to reach the result in a simple way and as soon as possible I have used the fact that the interactions of the chiral bosons are very similar to those of the W and W'. Therefore, I have just slightly corrected the PYTHIA code for the subprocess 142 (W' production) in the subroutines PYWIDT and PYRESD.

The subroutine PYWIDT calculates full and partial widths of resonances. As far as the decays of the gauge

and the chiral bosons are described by different formulas in the case of their decay into two massive fermions, I have substituted the matrix element

$$|\mathcal{M}_{W'}|^2 \;=\; 8M^2\left\{\left(g_V^2+g_A^2\right)\left[2-r_1-r_2-(r_1-r_2)^2\right]\right.$$
$$\left.+6\left(g_V^2-g_A^2\right)\sqrt{r_1r_2}\right\} \qquad (58)$$

in the expression for the W' width with the following expression

$$|\mathcal{M}_{\mathcal{U}}|^2 \;=\; 8M^2\left\{\left(g_V^2+g_A^2\right)\left[1+r_1+r_2-2\,(r_1-r_2)^2\right]\right.$$
$$\left.+6\left(g_V^2-g_A^2\right)\sqrt{r_1r_2}\right\} \qquad (59)$$

for the chiral bosons. Here $r_1=(m_1/M)^2$ and $r_2=(m_2/M)^2$ are squared ratios of the fermions masses m_1 and m_2 to the boson mass M.

The parameters PARU(131)/PARU(132) and PARU(133)/PARU(134) set the vector g_V/axial g_A couplings of the quark and lepton pairs to the heavy vector boson, correspondingly. These is a bug in the case of W' decay with non-equal vector and axial-vector couplings constants $|g_V|\neq|g_A|$, because the second line of eq. (58) is not coded into the program. Corrected in this way the subroutine should provide correct description of the W' decay with the W-like coupling constants PARU(131)=PARU(133)=1, PARU(132)=PARU(134)=-1 or the chiral $\mathcal{U}$ boson with following parameters

$$\mathrm{PARU}(131)=\frac{3}{2},\quad \mathrm{PARU}(132)=-\frac{1}{2},$$
$$\mathrm{PARU}(133)=\frac{1}{2},\quad \mathrm{PARU}(134)=-\frac{1}{2}. \qquad (60)$$

The subroutine PYRESD describes angular distributions of the resonances, which are completely different in the case of the gauge and chiral bosons. So, in accordance with the normalized distributions (47) and (48), the distribution

$$\mathrm{WT}=1+\mathrm{ASYM}\cdot\cos\hat\theta+\cos^2\hat\theta \qquad (61)$$

and its maximum value WTMAX=2+ABS(ASYM) for the gauge bosons, where the coefficient ASYM defines backward-forward asymmetry, should be replaced by simple formula

$$\mathrm{WT}=4\cos^2\hat\theta \qquad (62)$$

and a corresponding maximum value WTMAX=4 for the chiral bosons.

These changes open the possibility for full detector simulations of the production and diverse decay channels of the charged chiral bosons. The implementation of the neutral chiral bosons within PYTHIA cannot be fulfilled just as trivial modifications of already existing code for the Z' boson, because the code takes into account the full $\gamma^*/Z/Z'$ interference structure and is very complicated. In the same time the chiral bosons do not interfere with the gauge bosons and the new part of the code should be simplier and could be easily added by a PYTHIA expert.

I will continue with consideration of the golden discovery channels $pp\to\mathcal{U}/\mathcal{T}+X\to\ell\bar\nu+X'$ and $pp\to T^R/T^I+X\to\ell\bar\ell+X'$ for the chiral bosons at the LHC using CalcHEP package. Due to the enormous collider energy and the designed peak luminosity 10^{34} cm^{-2}s^{-1} they could be explored already at the very beginning, when the SM calibration processes $pp\to W+X\to\ell\bar\nu+X'$ and $pp\to Z+X\to\ell\bar\ell+X'$ are studied. So, the differential cross-sections for the gauge W (background) and the chiral $\mathcal{U},\mathcal{T}$ (signal) bosons as a function of the lepton transverse momentum are shown in the Fig. 6.

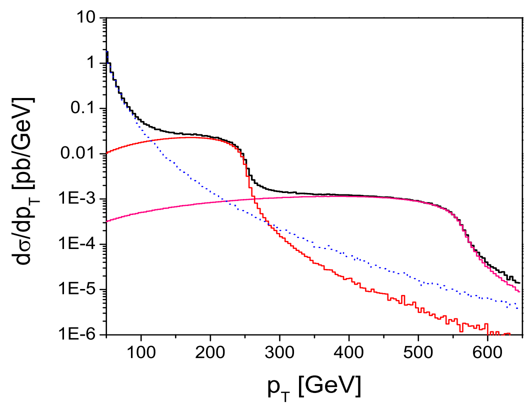

FIG. 6: The differential cross-sections for the gauge W (dotted) and the chiral $\mathcal{U},\mathcal{T}$ (solid) bosons as functions of the lepton transverse momentum.

The first shoulder corresponds to the resonance production of the light charged chiral bosons $\mathcal{U}^\pm$ with a differential cross-section of the order of 0.03 pb/GeV. It means, that after only an hour of data-taking at the peak luminosity, approximately 10 events should be within each bin with transverse lepton momentum around 200 GeV and 10 GeV bin's width. To see the second shoulder, corresponding to the heavy charged chiral boson, above one day of the data-taking is required.

However, the bad understanding of the detector resolution at the first runs and the uncertainties in the transverse momentum of the heavy bosons smear the Jacobian peak and the production of the chiral bosons could not be distinguished from the production of the gauge bosons. More data and detailed simulations will be needed. Therefore, the first crucial test of the discussed model at the LHC will be the observation of the peak at 1 TeV in the Drell–Yan dilepton channel (Fig. 7).

As it was discussed early, the shape of the peak as a function of the invariant dilepton mass is the same as for the gauge and chiral bosons. The only difference should

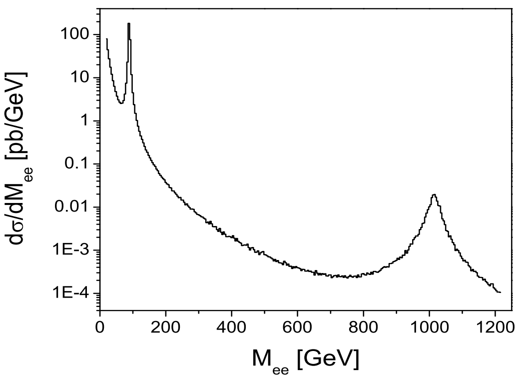

FIG. 7: The Z^0 and T^0 boson peaks in the differential cross-section as a function of the invariant dilepton mass.

be in the angular distributions of the lepton pairs, in particular, their p_T distributions. Instead the Jacobian peak at $p_T \simeq M_V/2$ for the gauge bosons, a wide bump, well below this point, is expected for the chiral boson in the lepton transverse distribution (see Fig. 2).

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[1] G. Arnison *et al.* (UA1 Collaboration), Phys. Lett. B **122**, 103 (1983); Phys. Lett.B **126**, 398 (1983).
[2] M. Banner *et al.* (UA2 Collaboration), Phys. Lett. B **122**, 476 (1983); P. Bagnaia *et al.* (UA2 Collaboration), Phys. Lett.B **129**, 130 (1983).
[3] P. Antilogus *et al.* (LEP Collaborations and LEP Electroweak Working Group), CERN-PPE/95-172, 1995.
[4] F. Abe *et al.* (CDF Collaboration), Phys. Rev. Lett. **74**, 2626 (1995).
[5] S. Abachi *et al.* (D0 Collaboration), Phys. Rev. Lett. **74**, 2632 (1995).
[6] Tevatron Electroweak Working Group, hep-ex/0703034.
[7] V.M. Abazov *et al.* (D0 Collaboration), Phys. Rev. Lett. **98**, 181802 (2007).
[8] CDF Collaboration, CDF Note 8665.
[9] V. D. Barger, R. J. N. Phillips, Collider Physics, Addison-Wesley publishing company, 1987.
[10] M. V. Chizhov, hep-ph/0609141.
[11] N. Kemmer, Proc. Roy. Soc. A **166**, 127 (1938).
[12] Supergravities in Diverse Dimensions, eds. A. Salam and E. Sezgin (North-Holland and World Scientific, 1989).
[13] A. Pukhov *et al.*, Preprint INP MSU 98-41/542, hep-ph/9908288; A. Pukhov, hep-ph/0412191.
[14] T. Sjöstrand, S. Mrenna, P. Skands, JHEP **05**, 026 (2006).
[15] V. K. Dobrev *et al.*, Bulg. J. Phys. **1**, 42 (1974).
[16] L. V. Avdeev, M. V. Chizhov, Phys. Part. Nucl. Lett. **2**, 7 (2005) [hep-th/9407067].
[17] L. V. Avdeev and M. V. Chizhov, Phys. Lett. B **321**, 212 (1994).
[18] M. V. Chizhov, Mod. Phys. Lett. A **8**, 2753 (1993) [hep-ph/0401217].
[19] M. V. Chizhov, hep-ph/9610220.
[20] H. Leutwyler, Phys. Lett. B **96**, 154 (1980).
[21] M. V. Chizhov, hep-ph/0008187.
[22] M. V. Chizhov, JETP Lett. **80**, 73 (2004) [hep-ph/0307100].
[23] O. Shankar, Nucl. Phys. B **204**, 375 (1982).
[24] M. B. Voloshin, Phys. Lett. B **283**, 120 (1992).
[25] V. N. Bolotov *et al.*, Phys. Lett. B **243**, 308 (1990); E. Frlež *et al.*. Phys. Rev. Lett. **93**, 181804 (2004).
[26] M. V. Chizhov, Phys. Part. Nucl. Lett. **2**, 193 (2005).
[27] D. Dubbers, talk at the Workshop on Precision Measurements at Low Energy, PSI, Villigen, Switzerland, 18-19 Jan 2007.
[28] W.-M. Yao *et al.*, Journal of Physics G **1**, 1 (2006).
[29] W. J. Marciano and A. Sirlin, Phys. Rev. Lett. **96**, 032002 (2006).
[30] M. V. Chizhov, hep-ph/0411098.
[31] M. V. Chizhov, hep-ph/0311360.
[32] S. S. Gershtein and Ya. B. Zeldovich, Zh. Eksp. Teor. Fiz. **29**, 698 (1955); R. P. Feynman and M. Gell-Mann, Phys. Rev. **109**, 193 (1958).
[33] M. Davier, hep-ph/0701163.
[34] G.W. Bennett *et al.*, Phys. Rev. Lett. **92**, 161802 (2004); Phys. Rev. D **73**, 072003 (2006).
[35] P. J. Mohr and B. N. Taylor, Rev. Mod. Phys. **77**, 1 (2005).
[36] B. Odom, D. Hanneke, B. D'Urso, and G. Gabrielse, Phys. Rev. Lett. **97**, 030801 (2006).
[37] G. Gabrielse, D. Hanneke, T. Kinoshita, M. Nio, and B. Odom, Phys. Rev. Lett. **97**, 030802 (2006).
[38] P. Cladé *et al.*, Phys. Rev. Lett. **96**, 033001 (2006).
[39] V. Gerginov *et al.*, Phys. Rev. A **73**, 032504 (2006).
[40] J. Pumplin *et al.*, JHEP **07**, 012 (2002).
[41] F. Abe *et al.*, Phys. Rev. Lett. **77**, 438 (1996).
[42] J. Huston *et al.*, Phys. Rev. Lett. **77**, 444 (1996); H. L. Lai *et al.*. Phys. Rev. D **55**, 1280 (1997).
[43] A. Abulencia *et al.* (CDF Collaboration), hep-ex/0611022.
[44] D0 Collaboration, D0 note 5191-CONF.
[45] G. Arnison *et al.* (UA1 Collaboration), Phys. Lett. B **129**, 273 (1983).
[46] CDF Collaboration, CDF Note 8747.
[47] V. M. Abazov *et al.* (D0 Collaboration), Phys. Lett. B **641**, 423 (2006).
[48] D0 Collaboration, D0 Note 5024-CONF.
[49] E. Boos *et al.*, hep-ph/0109068.
[50] C. Hof, T. Hebbeker, and K. Höpfner, CERN-CMS-NOTE-2006-117.

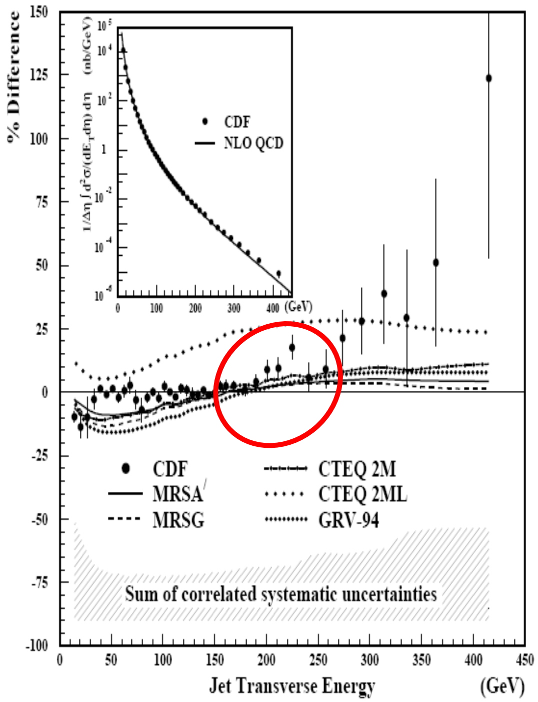
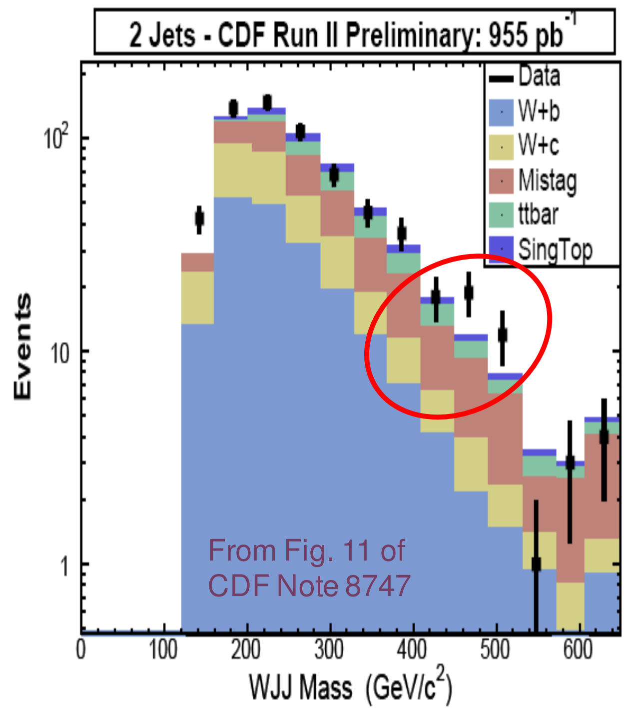
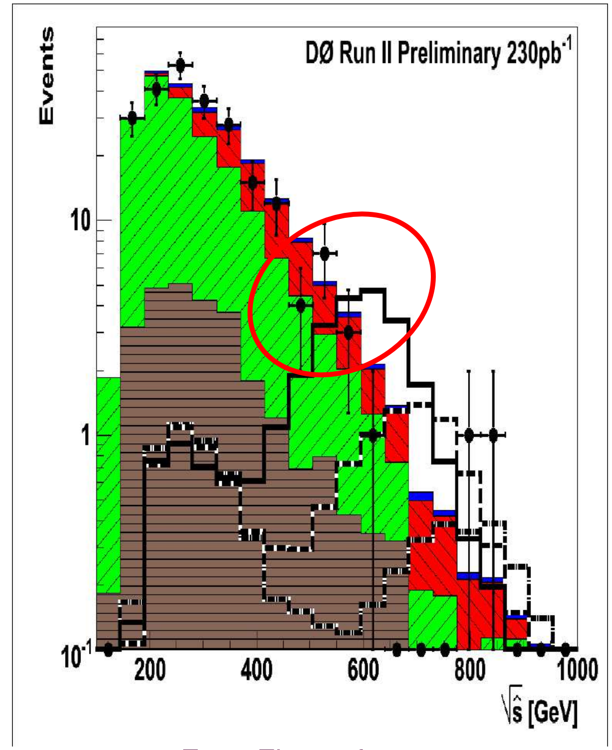


Fig. 1 of F. Abe *et al.*, PRL **77**(1996)438





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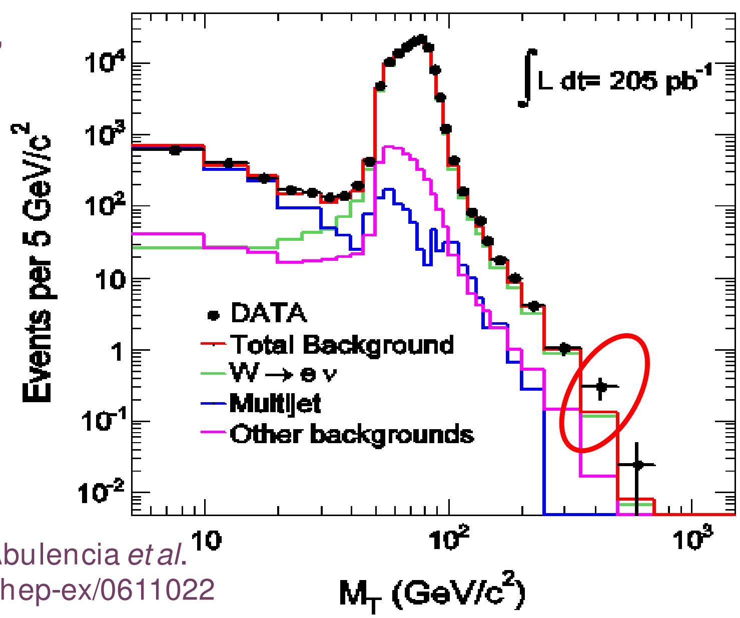

Fig. 1 of A. Abulencia *et al.* (CDF Coll.), hep-ex/0611022